

# Symmetry Breaking Patterns for the Little Higgs from Strong Dynamics

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We show how the symmetry breaking pattern of the simplest little Higgs model, and that of the smallest moose model that incorporates an approximate custodial SU(2), can be realized through the condensation of strongly coupled fermions. In each case a custodial SU(2) symmetry of the new strong dynamics limits the sizes of corrections to precision electroweak observables. In the case of the simplest little Higgs, there are no new light states beyond those present in the original model. However, our realization of the symmetry breaking pattern of the moose model predicts an additional scalar field with mass of order a TeV or higher that has exactly the same quantum numbers as the Standard Model Higgs and which decays primarily to third generation quarks.

## I. INTRODUCTION

In the Standard Model (SM) the Higgs mass parameter receives quadratically divergent radiative corrections from high scales, leading to a hierarchy problem. This suggests the existence of new physics at or close to a TeV that cancels these radiative corrections, thereby stabilizing the weak scale.

One interesting possibility, first explored in [1, 2, 3], is that the Higgs mass parameter is protected against radiative corrections because the Higgs is the pseudo Nambu-Goldstone boson (pNGB) of an approximate global symmetry. This idea has recently experienced a revival in the form of little Higgs theories [4, 5, 6, 7, 8], and also twin Higgs theories [9, 10]. These theories stabilize the weak scale against radiative corrections up to scales of order 10 TeV.

It is important to find ultra-violet (UV) completions of these theories that extend their validity to scales beyond 10 TeV. Non-supersymmetric weakly coupled UV completions of the simplest little Higgs [11] and simple group little Higgs [12] have been found. Supersymmetric UV completions of the little Higgs [13], (see also [14]), have also been constructed, in the context of the supersymmetric little hierarchy problem. There has also been work on the difficult problem of finding UV completions of these theories where the pattern of symmetry breaking is realized through strong dynamics [15, 16]. In this case the problem is complicated by the requirement that the new strong dynamics preserve a custodial SU(2) symmetry, so as not to generate large corrections to precision electroweak observables. An alternative approach has been to construct holographic little Higgs models in five dimensions [17], which are related to strongly coupled theories in four dimensions through the AdS/CFT correspondence.

In this paper we show that the characteristic symmetry breaking patterns of two well-motivated little Higgs theories can be realized through the condensation of strongly coupled fermions. We begin by dynamically realizing the symmetry breaking pattern of the simplest little Higgs model. Our construction does not require any additional light states beyond those of the original

model, and a custodial SU(2) symmetry of the new strong dynamics limits the size of corrections to precision electroweak observables. We then go on to consider the moose model of Chang and Wacker [8], which is the smallest extension of the minimal moose [5] which incorporates an approximate custodial SU(2) symmetry, and show that this symmetry breaking pattern can also be realized through strong dynamics. In the case of this ‘Next-to-Minimal Moose’ (NMM) model, our construction predicts the existence of an additional scalar field with mass of order a TeV that has exactly the same quantum numbers as the Standard Model Higgs, and which decays primarily to third generation quarks.

## II. THE SIMPLEST LITTLE HIGGS

Consider first the simplest little Higgs, which has the symmetry breaking pattern  $[SU(3) \times U(1) \rightarrow SU(2) \times U(1)]^2$ . The vector  $SU(3) \times U(1)$  subgroup, which is gauged, is broken down to the SM gauge group  $SU(2) \times U(1)$ , and 5 of the 10 Nambu-Goldstone bosons (NGBs) are eaten. The remaining 5 are actually pNGBs, and they constitute the Higgs of the SM and an additional SM singlet.

The interactions of the pNGBs are governed by universal low-energy theorems which are independent of any specific UV completion. This allows us to write down an effective field theory for the pNGBs valid at low energies, which takes the form of a non-linear sigma model.

We parameterize the pNGB degrees of freedom as  $h^a$  and  $\hat{h}^a$ . We write

$$\phi = \exp\left(\frac{i}{f} h^a t^a\right) \begin{pmatrix} 0 \\ 0 \end{pmatrix} \quad \hat{\phi} = \exp\left(\frac{i}{\hat{f}} \hat{h}^a t^a\right) \begin{pmatrix} 0 \\ 0 \end{pmatrix}, \quad (1)$$

where  $f$  and  $\hat{f}$  are the two independent breaking scales for each  $SU(3) \times U(1)$  global symmetry. The 5 matrices  $t^a$  span  $[SU(3) \times U(1)/SU(2) \times U(1)]$ :

$$\{t^a\} = \left\{ \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & i \\ 0 & -i & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & -1 \\ 0 & -1 & 0 \end{pmatrix}, \begin{pmatrix} 0 & 0 & i \\ 0 & 0 & 0 \\ -i & 0 & 0 \end{pmatrix}, \right.$$

$$\left\{ \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 0 \\ 1 & 0 & 0 \end{pmatrix}, \sqrt{\frac{2}{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -1 \end{pmatrix} \right\}.$$

In general the low energy effective Lagrangian for the fields  $h^a, \hat{h}^a$  will contain all operators involving  $\phi$  and  $\hat{\phi}$  consistent with the non-linearly realized  $SU(3) \times U(1)$  symmetry, suppressed by the cutoff of the non-linear sigma model, which we denote by  $\Lambda$ . The scale  $\Lambda$  is constrained to be less than or of order  $4\pi f$ , where the upper bound is at strong coupling. The coefficients of these operators are determined by the UV physics.

One of the operators allowed by symmetry takes the form

$$c \frac{|\phi^\dagger D_\mu \phi|^2}{\Lambda^2}, \quad (3)$$

where the value of the coefficient  $c$  depends on the specific UV completion. This operator violates the approximate custodial  $SU(2)$  symmetry of the Standard Model. Its effect is to alter the ratio of the masses of the W and Z gauge bosons from the Standard Model prediction, and it is therefore very tightly constrained by experiment. A potential problem with any UV completion where this pattern of symmetry breaking is realized through strong dynamics is that  $c$  is expected to be of order  $(4\pi)^2$ , constraining the cutoff  $\Lambda$  to be of order 50 TeV or higher. Since little Higgs theories can only stabilize the weak scale up to energies of order 10 TeV or so, this reintroduces fine-tuning.

Is there a way that this operator can be suppressed? A linear realization of this symmetry breaking pattern is instructive: Consider two scalar fields  $\Phi$  and  $\hat{\Phi}$  which transform as a  $3_{-\frac{1}{3}}$  under the  $SU(3) \times U(1)$  gauge symmetry. The Lagrangian is

$$(|D_\mu \Phi|^2 + m^2 |\Phi|^2 - \lambda |\Phi|^4) + (\Phi \rightarrow \hat{\Phi}). \quad (4)$$

After symmetry breaking one linear combination of the NGBs from  $\Phi$  and  $\hat{\Phi}$  is eaten, while the orthogonal linear combination which contains the SM Higgs survives. We denote the VEV of  $\Phi$  by  $f$  and that of  $\hat{\Phi}$  by  $\hat{f}$ . The pNGB fields of the non-linear model are those degrees of freedom which survive after integrating out the radial modes of the fields  $\Phi$  and  $\hat{\Phi}$  in the linear model.

The key observation is that in the limit that the gauge symmetry is turned off this potential, Eq. (4), has an accidental  $SO(6)^2$  global symmetry which is broken to  $SO(5)^2$ , and the 10 NGBs can just as well be thought of as arising from this symmetry breaking pattern. Once we gauge the vector  $SU(3) \times U(1)$  subgroup of the global symmetry again, 5 of these NGBs are eaten, while the remaining 5 survive in the low energy theory as pNGBs.

Consider now the situation where the breaking pattern  $SO(6)^2$  to  $SO(5)^2$  is realized non-linearly. Since the symmetry breaking pattern now preserves the custodial  $SU(2)$  of the SM, operators such as Eq. (3) are forbidden to leading order. They are only generated through loops involving those interactions (gauge and Yukawa

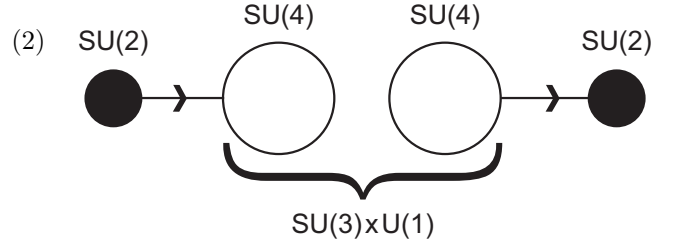


FIG. 1: A UV completion of the simplest little Higgs. The flavor  $SU(4)^2$  symmetry has gauged subgroup  $SU(3) \times U(1)$ . The curly braces indicate that only the diagonal subgroup of  $SU(4)^2$  is gauged. The black filled-in circle is a condensing  $SU(2)$  group. Each link field is a collection of four  $SU(2)$  fundamentals that transforms like a fundamental (anti-fundamental) of  $SU(4)$ .

couplings) that violate the custodial  $SU(2)$  symmetry, and therefore the coefficient  $c$  is at most of order one. This allows the scale  $\Lambda$  in Eq. (3) to be as low as 5 TeV without conflicting with the constraints from precision electroweak measurements.

We see from this that the first step to finding a strongly coupled UV completion of the simplest little Higgs is to find a way to break  $SO(6)^2$  to  $SO(5)^2$  through strong dynamics. We now explain how this can be accomplished. Notice that  $SU(4)$  is the double-cover of  $SO(6)$ , and  $Sp(4)$  is the double-cover of  $SO(5)$ ; in both cases the Lie algebras are isomorphic. The problem is therefore to break  $SU(4)^2$  to  $Sp(4)^2$  through strong dynamics. Consider an  $SU(2)$  gauge theory with 4 fields  $\psi_{\alpha i}$  in the fundamental representation. Here  $\alpha$  is an  $SU(2)$  gauge index and  $i$  takes values 1 through 4. When the  $SU(2)$  gauge theory gets strong we expect a condensate  $\langle \epsilon_{\alpha\beta} \psi_{\alpha i} \psi_{\beta j} \rangle \propto J_{ij}$  to form along the gauge singlet direction. This is antisymmetric in the indices  $i$  and  $j$ , thereby breaking the  $SU(4)$  global symmetry down to  $Sp(4)$ .<sup>\*</sup> In order to break  $SU(4)^2$  to  $Sp(4)^2$  we merely begin with two copies of the  $SU(2)$  gauge theory—each with 4 fields in the fundamental representation. We label the two copies of the  $\psi$ 's as  $\psi_{\alpha i}$  and  $\hat{\psi}_{\alpha i}$ , and the corresponding condensates  $J_{ij}$  and  $\hat{J}_{ij}$ .

We are free to weakly gauge a vector  $SU(3) \times U(1)$  subgroup of the  $SU(4)^2$  global symmetry, as shown in Figure 1. Without loss of generality we may take the indices  $i = 1, 2, 3$  to be  $SU(3)$  gauge indices, while  $i = 4$  is not an  $SU(3)$  gauge index. The  $SU(3)$  gauge symmetry is anomaly free provided the fields in  $\psi$  transform in the fundamental representation and the fields in  $\hat{\psi}$  in the anti-fundamental. We choose the charges of the  $U(1)$  such that the fields in  $\psi$  transforming as the fundamental of  $SU(3)$  have charge  $+1/6$ , while the  $SU(3)$  singlet has

<sup>\*</sup> More generally, any condensing  $Sp(2N)$  with four fundamental flavors of fermions can be used to break  $SU(4)$  to  $Sp(4)$  [3].

charge  $-1/2$ . The anti-fundamental of  $SU(3)$  in  $\hat{\psi}$  has  $U(1)$  charge  $-1/6$  and the  $SU(3)$  singlet charge  $+1/2$ . After condensation this  $SU(3) \times U(1)$  gauge symmetry is broken to the  $SU(2) \times U(1)$  of the SM.

The low energy effective theory can once again be described by a non-linear sigma model. The matrix  $J_{ij}$  has the structure

$$J = \begin{pmatrix} i\sigma^2 & 0 \\ 0 & i\sigma^2 \end{pmatrix}. \quad (5)$$

The matrix  $J$  is left invariant under  $Sp(4)$  rotations, but all other generators of  $SU(4)$  are broken. We parameterize the five resulting low energy degrees of freedom as  $\pi^a$ . Define

$$A = f \exp\left(\frac{i}{f} \pi^a T^a\right) J \exp\left(\frac{i}{f} \pi^a T^a\right)^T. \quad (6)$$

Here the 5 matrices  $T^a$  are the generators of  $SU(4)/Sp(4)$ .

$$\{T^a\} = \left\{ \begin{pmatrix} 0 & i\sigma^a \\ -i\sigma^a & 0 \end{pmatrix}, \begin{pmatrix} 0 & I \\ I & 0 \end{pmatrix}, \begin{pmatrix} I & 0 \\ 0 & -I \end{pmatrix} \right\}. \quad (7)$$

The matrix  $\hat{J}_{ij}$  has exactly the same form as the matrix  $J_{ij}$ , and we parameterize the corresponding low energy degrees of freedom by  $\hat{\pi}^a$ , collected in a matrix  $\hat{A}$ .

By choosing to keep only the  $SU(3) \times U(1)$  subgroup of each  $SU(4)$  symmetry manifest, we can immediately carry over all results of the original simplest little Higgs model to this new construction. To see this explicitly, consider the decomposition of  $SU(4)$  into its  $SU(3) \times U(1)$  subgroup. We can identify the  $\pi^i$ 's and  $\hat{\pi}^i$ 's with the  $h^i$ 's and  $\hat{h}^i$ 's up to constants of proportionality, since they have exactly the same transformation properties under the nonlinearly realized  $SU(3) \times U(1)$ :

$$A = \begin{pmatrix} 0 & \phi_3^* & -\phi_2^* & \phi_1 \\ -\phi_3^* & 0 & \phi_1^* & \phi_2 \\ \phi_2^* & -\phi_1^* & 0 & \phi_3 \\ -\phi_1 & -\phi_2 & -\phi_3 & 0 \end{pmatrix}. \quad (8)$$

In the absence of any explicit symmetry breaking, the low energy effective Lagrangian for the pNGBs consists of all operators involving  $A$  and  $\hat{A}$  consistent with the non-linearly realized  $SU(4)^2$  symmetry. Equivalently, the low energy effective Lagrangian consists of all possible operators involving  $\phi$  and  $\hat{\phi}$  consistent with the nonlinearly realized  $[SU(3) \times U(1)]^2$  symmetry, but with additional relations among the coefficients of the various terms enforced by the larger  $SU(4)^2$  global symmetry. In particular, dangerous operators which violate the custodial  $SU(2)$  symmetry, such as that in Eq. (3), are forbidden at leading order.

Any potential for the pNGBs can arise only from those interactions that explicitly violate the  $SU(4)^2$  global symmetry, in particular the  $SU(3) \times U(1)$  gauge symmetry and the Yukawa couplings. In the low-energy effective Lagrangian these interactions can be written down explicitly in terms of the fields  $\phi$  and  $\hat{\phi}$ , exactly as

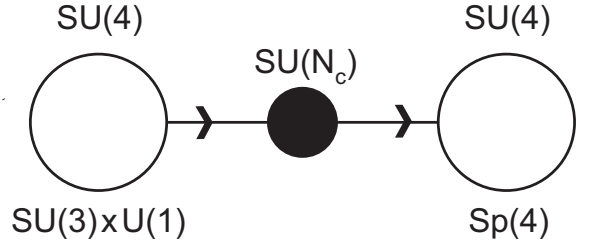


FIG. 2: An alternate moose for the  $SO(6) \rightarrow SO(5)$  symmetry breaking pattern. The moose has an  $SU(4)^2$  flavor symmetry, with gauged subgroup  $Sp(4) \times SU(3) \times U(1)$ . One condensing  $SU(N_c)$  group is also present, and each link field is a bifundamental fermion under the neighboring global symmetries.

in the original model. The SM fermions can be embedded, anomaly-free, under  $SU(3) \times U(1)$ , by promoting SM doublets of  $SU(2)$  to triplets or anti-triplets as required [7],[18]. The Yukawa interactions of the low energy effective theory arise from non-renormalizable interactions between the SM fermions and the fermions that condense to form the Higgs. Consider, for concreteness, the top Yukawa coupling. The third-generation left-handed quark doublet,  $q$ , is enlarged to an  $SU(3)$  triplet with  $U(1)$  charge  $1/3$ ,  $Q = (q, U)$ . This triplet couples to two  $SU(3)$  singlets  $U_1^c$  and  $U_2^c$  through the following terms:

$$\frac{\lambda_1}{4\pi f^2} (\psi_i \psi_4)^\dagger Q_i U_1^c + \frac{\lambda_2}{4\pi f^2} (\hat{\psi}^i \hat{\psi}^4) Q_i U_2^c, \quad (9)$$

where only the  $SU(3)$  gauge indices are shown explicitly. This leads to the couplings below in the low energy effective theory after the  $SU(2)$  groups get strong:

$$y_1 \phi_i Q_i U_1^c + y_2 \hat{\phi}_i Q_i U_2^c. \quad (10)$$

These interactions, which are familiar from the original model, give rise to the top Yukawa coupling while eliminating the one-loop quadratic divergences from the top loop. We leave the important question of the UV origin of the higher dimensional operators in Eq.(9) for future work.

It is interesting to note that there is another configuration which also generates exactly the symmetry breaking pattern of the simplest little Higgs with a custodial  $SU(2)$  symmetry. As pointed out in [16], the global symmetry breaking pattern  $G \rightarrow H$  with a subgroup  $F$  of  $G$  gauged has the same low energy dynamics as a two-site nonlinear sigma model with global symmetry breaking pattern  $G^2 \rightarrow G$  and gauged subgroup  $H \times F$ , in the limit that the gauge coupling constant of  $H$  is large. This two-site model can in principle be UV completed, provided the  $G^2 \rightarrow G$  breaking pattern is a generalization of QCD-type confining dynamics.

Since the pattern  $SO(6) \rightarrow SO(5) \simeq SU(4) \rightarrow Sp(4)$ , we can use this technique to realize symmetry breaking in the simplest little Higgs. The appropriate construction is shown in Figure (2); to replicate exactly the simplest

little Higgs we simply repeat this construction twice and gauge the same  $SU(3) \times U(1)$  symmetry in each case.

### III. THE NEXT-TO-MINIMAL MOOSE

The NMM model [8] is built around the symmetry breaking pattern  $SO(5)^2 \rightarrow SO(5)$ . The same pattern also arises in several other little Higgs models [19]. We now show that the symmetry breaking pattern  $Sp(4)^2 \rightarrow Sp(4)$ , which is equivalent to  $SO(5)^2 \rightarrow SO(5)$ , can be realized through strong dynamics. Consider an  $SU(N_c)$  gauge group, with a set of four fermions,  $\chi_{\alpha i}$ , in the fundamental representation. Here  $\alpha$  represents an  $SU(N_c)$  gauge index and  $i$  labels the fermions from 1 through 4. We also add a set of four fermions in the conjugate representation  $\hat{\chi}_i^\alpha$ . When the  $SU(N_c)$  theory gets strong, a condensate  $\langle \hat{\chi}_i^\alpha \chi_{\alpha j} \rangle \propto \delta_{ij}$  forms and breaks the  $SU(4)_L \times SU(4)_R$  flavor symmetry to the diagonal  $SU(4)$ . We label the 15 resulting NGBs that are produced by  $\pi^a$ , and define

$$X = f \exp(2i\pi^a T^a / f), \quad (11)$$

where the matrices  $T^a$  are generators of  $SU(4)$ .

We also add to the theory a non-renormalizable term

$$\frac{m^2}{(4\pi f^2)^2} \text{Tr} \left[ (\chi \hat{\chi}) J (\chi \hat{\chi})^T J \right] \sim m^2 \text{Tr} [X J X^T J] \quad (12)$$

which is allowed by the gauge symmetries. The effect of this term is to explicitly break the global  $SU(4)^2$  symmetry to  $Sp(4)^2$ , thereby giving a mass of order  $m$  to 5 of the 15 NGBs. With the addition of this term the pattern of global symmetry breaking is in fact  $Sp(4)^2 \rightarrow Sp(4)$ , which accounts for the 10 surviving NGBs. The unbroken global symmetry, the diagonal  $Sp(4)$ , contains the custodial  $SU(2)$  symmetry we desire.

To reproduce the NMM gauge symmetry breaking pattern, we weakly gauge an  $SU(2)_L \times SU(2)_{L'}$  subgroup of the  $SU(4)_L$  symmetry, and an  $SU(2)_R \times U(1)_{R'}$  symmetry of the  $SU(4)_R$  global symmetry, as shown in Figure 3. Here  $U(1)_{R'}$  is the diagonal generator of the other  $SU(2)_{R'}$  subgroup contained in  $SU(4)_R$ . The gauge symmetry is broken to the diagonal  $SU(2) \times U(1)$  of the SM and 6 NGBs are eaten. The remaining 4 light pNGBs constitute a single complex Higgs doublet. The custodial  $SU(2)$  symmetry is preserved by the kinetic terms up to small corrections arising from integrating out the states of mass  $m$ . Provided  $m$  is larger than or of order a TeV, these are under control. This is in contrast to the original minimal moose where corrections to the precision electroweak observables arising from the kinetic terms are in general very large [20].

It is now straightforward to reproduce the complete symmetry breaking pattern of the NMM model by repeating the breaking pattern  $SO(5)^2 \rightarrow SO(5)$  multiple times. Here we limit ourselves to constructing a model

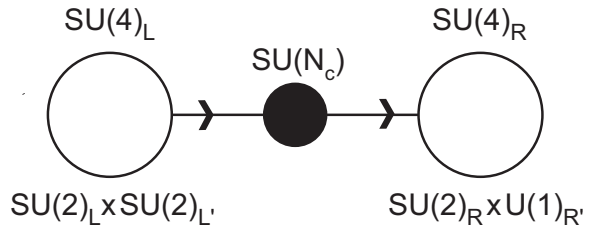


FIG. 3: A UV completion for the minimal moose. The theory has a global  $SU(4)_L \times SU(4)_R$  flavor symmetry with the indicated (written along the bottom) gauged subgroup. There is also a condensing  $SU(N_c)$  gauge group. Each link represents a bifundamental fermion under the indicated symmetries.

with just a single link that reproduces the main features of the NMM model.

The fermion sector is straightforward and anomaly-free. Each SM  $SU(2)$  doublet,  $q$  or  $l$ , transforms under  $SU(2)_L$  and emerges from a fundamental of  $SU(4)_L$ , which we denote by  $\hat{q}$  or  $\hat{l}$ . A second doublet,  $q'$  or  $l'$ , which transforms under  $SU(2)_{L'}$  fills out  $\hat{q}$  or  $\hat{l}$ . The fields  $\hat{q}$  and  $\hat{l}$  carry  $U(1)_{R'}$  charges equal to the SM hypercharges of  $q$  and  $l$  respectively. In addition, corresponding to each  $q'$  or  $l'$  is a field  $q'^c$  or  $l'^c$  that is vector-like with respect to it. The SM  $SU(2)$  singlet fields,  $U^c, D^c$  and  $E^c$ , only carry  $U(1)_{R'}$  gauge quantum numbers with charges again equal the corresponding SM hypercharges. The quark Yukawa couplings emerge from

$$\mathcal{L} \supset \frac{1}{4\pi f^2} \begin{pmatrix} q \\ q' \end{pmatrix} \cdot (\chi \hat{\chi}) \begin{pmatrix} 0 \\ 0 \\ \lambda_U U^c \\ \lambda_D D^c \end{pmatrix} + \lambda' f q' q'^c. \quad (13)$$

One linear combination of  $U^c$  and  $U'^c \subset q'^c$  acquires a mass of order  $f$ , while the other linear combination is the usual SM singlet,  $u^c$ . A similar story holds for the down-type quarks. The invariance of the first term under  $SU(4)_L$  guarantees the cancellation of one loop quadratic divergences from the quark sector. The generalization to leptons is straightforward.

The states with mass  $m$  are not present in the original NMM model, but constitute a firm prediction of our construction. In particular, they will be present in any UV completion that realizes the pattern of symmetry breaking in the same manner. These states are composed of a field with the same SM quantum numbers as the Higgs that decays predominantly into third generation quarks, and a SM singlet.

In conclusion, we have shown how to obtain the symmetry breaking patterns of the simplest little Higgs and the smallest moose with an approximate custodial  $SU(2)$  symmetry from strong dynamics. Our hope is that this will open the door to the construction of completely realistic, UV completions of the little Higgs from strong dynamics.

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